

METHODS OF SUPERSYMMETRY BREAKING

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Abstract

This is a review of basic ideas and mechanisms encountered in the supersymmetry breaking problem at the global level, in supergravity models, and in superstring theory.

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1. Introduction

This conference is a reunion of true SUSY believers, so there is no need to argue that supersymmetry is really a symmetry of particle physics. It is clear that we are all facing here a long overdue problem why SUSY has not been seen at low energies. If it is a “good”, exact symmetry, it must be realised in a spontaneously broken mode, because only in this case can we use it to make definite predictions for superparticle masses and couplings. This is assuming that we understand the origin of its breaking – the super-Higgs mechanism. Unfortunately, this part of the supersymmetric standard model is still missing, which explains the rather academic title of this talk; it is intended as an introduction to the basic ideas and mechanisms of supersymmetry breaking.

In the standard model, electroweak symmetry is broken by a non-zero vacuum expectation value (VEV) of the Higgs doublet. In the case of supersymmetry, the analogous order parameters are the VEVs of auxiliary fields belonging to either chiral or vector multiplets. As explained in standard textbooks [1], auxiliary fields are introduced in order to close the off-shell supersymmetry algebra. Under a supersymmetry transformation, the fermion ψ belonging to a chiral multiplet transforms as

$$\delta_\epsilon \psi = i\sigma^m \bar{\epsilon} \partial_m A + \epsilon F \quad (1)$$

where A is the scalar partner of ψ and F is the auxiliary component of the multiplet. The latter does not contain any physical degree of freedom. After using lagrangian field equation, it becomes a function of physical fields: $F = F(A, \dots)$. The VEV of F is determined by further use of equations of motion, including minimisation of the scalar potential *etc.* If it turns out to be non-zero,

$$\langle 0|F|0\rangle = F(\langle 0|A|0\rangle, \dots) \equiv M_S^2, \quad (2)$$

supersymmetry is broken spontaneously at mass scale M_S . This is easy to understand. By looking at eq.(1) we see that

$$\langle 0|[\epsilon Q, \psi]|0\rangle = \langle 0|\epsilon F|0\rangle, \quad (3)$$

where Q is the supercharge operator. Then

$$\langle 0|\epsilon F|0\rangle \neq 0 \Rightarrow \epsilon Q|0\rangle \neq 0 \Rightarrow e^{i\epsilon Q}|0\rangle \neq |0\rangle, \quad (4)$$

so the vacuum state which carries non-zero supercharge is *not* invariant under supersymmetry transformations. Furthermore, it can be shown that a massless fermion – the goldstino – must be present in the spectrum, populating degenerate states obtained from the vacuum by SUSY transformations. In the case of F -type breaking this is exactly the fermion ψ which transforms under (1) into the auxiliary field acquiring a non-zero VEV. Another type of SUSY breaking, the so-called D -type breaking may occur in the presence of vector multiplets. A vector multiplet contains a gauge boson and a gaugino λ which transforms as

$$\delta_\epsilon \lambda = \sigma^{mn} \epsilon F_{mn} + i\epsilon D, \quad (5)$$

where F_{mn} is the gauge field strength and D is the auxiliary component of the multiplet. By an argument similar to (4), a non-zero VEV of D breaks SUSY, with the gaugino identified as the goldstino. If supersymmetry is gauged, *i.e.* promoted to a local symmetry, then the goldstino degrees of freedom are absorbed by the massive spin $\frac{3}{2}$ gravitino as its helicity $\pm\frac{1}{2}$ components.

The computation of F and D VEVs is a dynamical problem. It may be simple in the case of weakly interacting globally supersymmetric theories and supergravity, and possibly more difficult in the presence of strong interactions, but the basic idea is always the same: use field equations to determine auxiliary VEVs. The form of field equations depends on a particular

model. The universal feature is the necessary presence of massless goldstinos in spontaneously broken SUSY models. This provides an intuitive criterion for SUSY breaking: the breaking can occur only if there is a massless fermion in the spectrum – a potential goldstino. The most sophisticated and rigorous version of this argument is called the Witten index theorem [2]. I will discuss separately the cases of global SUSY, supergravity and superstring theory.

2. Globally Supersymmetric Renormalisable QFT

A globally supersymmetric QFT is completely specified by the superpotential $W(\Phi)$, an analytic function of chiral superfields Φ . The requirement of renormalisability restricts $W(\Phi)$ to a polynomial of degree 3 in Φ 's. The classical equations of motion for the auxiliary fields are

$$\bar{F}_{\bar{\Phi}} = \frac{\partial W}{\partial \Phi} \Big|_{\Phi=A} , \quad D_a = g_a \sum_A A^\dagger T^a A + \xi^a , \quad (6)$$

where g_a and T^a are the gauge group couplings and generators, respectively, and ξ^a is the Fayet-Iliopoulos parameter that may be non-zero for an index a associated to a $U(1)$ subgroup only. The classical scalar potential is

$$V(A) = \sum_{\Phi} |F_{\Phi}|^2 + \sum_a D_a^2 \geq 0 , \quad (7)$$

with the auxiliary fields given by eq.(6). In the weak coupling limit, the potential can be minimised to determine all VEVs and to see whether supersymmetry is broken or not. For instance, showing that $V > 0$ in the vacuum is completely sufficient to prove that some auxiliary fields acquire non-zero VEVs and hence SUSY is broken. This procedure can be a posteriori justified if it happens that all fields are weakly interacting at the SUSY breaking scale. As usual, life is not so simple: it turns out that supersymmetry remains unbroken in the minimal extension of the supersymmetric standard model. A completely new, “hidden” sector is necessary to trigger SUSY breaking. For electroweak symmetry it was sufficient to introduce one Higgs doublet with a simple potential, whereas in the case of supersymmetry one needs at least several chiral multiplets with a complicated superpotential and/or Fayet-Iliopoulos terms associated with exotic $U(1)$'s, each of them bringing in one more vector supermultiplet. In these types of models, the supersymmetry breaking scale M_S is introduced by hand. Another possibility is non-perturbative supersymmetry breaking due to condensates, *i.e.* non-zero VEVs of composite fields playing the role of auxiliary components [2, 3, 4]. M_S can then be determined from the strong interaction scale of “supercolour” forces that cause dynamical supersymmetry breaking, which may seem to be more natural than putting it in by hand. Supercolour theories are not too difficult to construct; an important ingredient is the absence of the mass gap, allowing the existence of composite goldstinos. The main problem, however, common to weakly and strongly coupled hidden sectors, is how to communicate SUSY breaking to the observable sector of quarks, squarks *etc.* A complicated system of “messengers” [5] must be employed in order to generate squark and gaugino masses. The main virtue of this approach, advertised by its proponents, is that the physics is fully contained below 1 TeV, staying away from the traps and *zasadzkas* of quantum gravity, strings *etc.* In principle, this is a completely calculable scheme, but in practice all viable models are very complicated and involve a great deal of theoretical uncertainty.

3. Local Supersymmetry and Standard Supergravity

As a consequence of the supersymmetry algebra which includes also the momentum operator, gauging supersymmetry automatically brings into the game gravity and the associated

parameters – the Planck mass $M_P \sim 10^{19}$ GeV and the coupling $\kappa \sim 1/M_P$. The gravitino $\psi_{3/2}$ is the spin $\frac{3}{2}$ gauge fermion of supersymmetry which belongs to the gravitational supermultiplet together with the spin 2 graviton. All known forces can be unified in the framework of supergravity. The theory is no longer renormalisable, but as far as SUSY breaking is concerned, the lack of renormalisability can be turned into advantage: higher dimensional interactions provide a natural “messenger” system for communicating SUSY breaking to the observable sector. Assuming that the scale of SUSY breaking VEVs in the hidden sector is of order Λ , and that higher-dimensional interactions $\mathcal{O}(\kappa^2)$ are responsible for the super-Higgs effects, we have $m_{3/2} \sim \Lambda^3/M_P^2$. A gravitino mass of order of 1 TeV can be then generated by the hidden sector dynamics at $\Lambda \sim 10^{13}$ GeV.

Together with relaxing the renormalisability requirement, there comes a possible field-dependence of parameters which are constrained to be constant in the global case. It is encoded in the Kähler potential K , in the superpotential W , and in the gauge functions f_a , which depend on chiral superfields [1]. W and f_a ’s are analytic while K is real. The chiral superfields (and the corresponding scalars), generically denoted by A , will be divided into the observable ones - q , and the hidden ones - ϕ . In order to analyse SUSY breaking by hidden VEVs, it is convenient to measure them in M_P units, which can be done by a simple rescaling that renders all ϕ ’s dimensionless. The field dependence can be seen in the following formulas for the wave-function factors \mathcal{Z} , Yukawa couplings Y , and the gauge couplings g_a :

$$\mathcal{Z}_{IJ}(A) = \frac{\partial^2 K(\phi, q)}{\partial q^I \partial \bar{q}^{\bar{J}}} , \quad Y_{IJK}(A) = \frac{\partial^3 W(\phi, q)}{\partial q^I \partial q^J \partial q^K} , \quad \frac{1}{g_a^2(A)} = \text{Re} f_a(\phi, q) . \quad (8)$$

Since one is mostly interested in the VEVs of hidden fields, expected to be much bigger than the observable VEVs, one can expand in powers of q ’s:

$$K = \kappa^{-2} \widehat{K}(\phi) + \mathcal{Z}_{IJ}(\phi) q^I \bar{q}^{\bar{J}} + \dots , \quad (9)$$

$$W = \widehat{W}(\phi) + Y_{IJK}(\phi) q^I q^J q^K + \dots \quad (10)$$

Note that the hidden Kähler potential \widehat{K} is dimensionless while the superpotential \widehat{W} has mass dimension 3, therefore its size is set by the scale Λ *i.e.* $\widehat{W} \sim \Lambda^3$. The supergravity version of the auxiliary field equations (6) is

$$\bar{F}_{\bar{\phi}} = \kappa^2 e^{\widehat{K}/2} \left(\frac{\partial^2 \widehat{K}}{\partial \phi \partial \bar{\phi}} \right)^{-1} \left(\frac{\partial \widehat{W}}{\partial \phi} + \widehat{W} \frac{\partial \widehat{K}}{\partial \phi} \right) + \dots \quad (11)$$

The D components are also given by expressions similar to eq.(6), however since F -type breaking is very easy to achieve, there is really no need to consider D -type breaking.

The formula for the scalar potential is slightly more complicated than eq.(7), and there is one important difference: it is not positive definite. There is also another difference: in order to find the vacuum it is no longer sufficient to minimise this potential. The gravitational equations of motion, which in the supergravity case play the role of gauge field equations, are equally important. If the minimum of the potential occurs at non-zero vacuum energy, the gravitational background has a non-zero curvature. A flat Minkowski background requires $V = 0$ at the minimum,^a which unlike in the global case, turns out to be compatible with broken SUSY. After ensuring that the classical minimum occurs at $V = 0$ at the classical level, it is not clear what to do with quantum corrections. Because of this, the famous cosmological constant inevitably gets in the way. There is no room for a separate “adjusting” of the cosmological constant without

^aUnless one considers more complicated, cosmological solutions with space-time dependent scalar fields.

ruining the mass relations *etc.* that follow from spontaneously broken supersymmetry. A possible procedure is to construct a model with a vanishing tree-level cosmological constant, derive the spectrum, couplings *etc.* and then analyse quantum corrections assuming the existence of a physical ultraviolet cutoff [6, 7].

Non-vanishing VEVs of hidden auxiliary components trigger spontaneous supersymmetry breaking, generating the gravitino mass

$$m_{3/2} = \left(\frac{1}{3} \frac{\partial^2 \widehat{K}}{\partial \phi \partial \bar{\phi}} F_{\phi} \bar{F}_{\bar{\phi}} \right)^{1/2} = \kappa^2 \langle 0 | e^{\widehat{K}(\phi)/2} \widehat{W}(\phi) | 0 \rangle . \quad (12)$$

It is important to be aware that the second part of this equation, familiar to model-builders, is correct under the assumption of $V = 0$ at the minimum, it is therefore sensitive to a possible fine-tuning of the cosmological constant. Note that a hidden superpotential $\widehat{W} \sim \Lambda^3$ does indeed generate $m_{3/2} \sim \Lambda^3/M_P^2$. As a result of higher-dimensional interactions between the observable and hidden sectors implied by the underlying supersymmetry, the observable scalars acquire masses $\mathcal{O}(m_{3/2})$. The exact expressions for these masses depend on details of the Kähler potential, *e.g.* $\mathcal{Z}_{I\bar{J}}$ factors *etc.*, therefore there is no reason to expect any special mass pattern.

When it comes to actual model building, there is no problem with constructing SUSY-breaking hidden sector superpotentials [8]. This can be achieved even by one chiral multiplet with a linear superpotential, like in the Polonyi model. In this case the hidden scale, hence effectively M_S , is introduced by hand. A more “natural” scenario is offered by no-scale models [9], where the Kähler potential is adjusted in such a way that a constant superpotential breaks supersymmetry with an identically vanishing scalar potential at the tree level. M_S is determined then by radiative corrections to be of the same order as the electroweak scale. The bottom line is a softly-broken supersymmetric low-energy effective field theory obtained from the supergravity lagrangian by taking the limit $\kappa \rightarrow 0$ while keeping $m_{3/2}$ fixed [10]. SUSY breaking can then be parametrised by a finite number of parameters. In the simplest supergravity models there are at least five such parameters: universal scalar mass m_0 , gaugino mass $m_{1/2}$, higgsino mass parameter μ , and two parameters, A and B , which specify the scalar potential. It is clear however that in the absence of renormalisability there is no rigorous guiding principle for selecting one hidden sector or another, therefore it is not possible to make a definite prediction for the structure of soft-breaking terms.

To summarise, supergravity provides a natural setting for SUSY breaking and a messenger system for feeding this breaking down to the supersymmetric standard model sector. On the other hand, the lack of renormalisability and the cosmological constant problem do clearly reduce its predictive power. First of all, supergravity by itself gives no indication about details of hidden sectors that are necessary to derive the properties of low-energy softly broken theory. Furthermore, even if one starts from a definite model at the classical level, it is not clear whether a consistent treatment is possible for quantum corrections [7]. Certainly, an ultraviolet cutoff is necessary in order to study the stability of the $M_P - M_S$ hierarchy and other phenomenological problems.

4. Superstring Theory

There is only one or at worst a few superstring theories^b – heterotic, type I, II *etc.* – but there are millions of four-dimensional models corresponding to apparently degenerate ground states of the same theory. The present understanding of short-distance superstring dynamics is not sufficient to select one particular model, or a class of models, so it is better to pursue a general analysis. Each particular model contains one parameter, the string mass

^b The reason for this hesitation should become clear at the end of the talk.

scale $M \sim M_P$, and its low-energy limit is described by a supergravity theory with definite K , W and f 's. The physical parameters like masses and couplings depend then on the VEVs of hidden and observable scalar fields. As far as SUSY breaking is concerned, the fundamental problem is to understand how $M_S \ll M$ can be generated by these VEVs. The breaking may involve effects associated with the extended string nature or it may be simply a field-theoretical phenomenon.

4.1. Stringy SUSY Breaking: Twisted and Magnetised Tori

Many four-dimensional superstring models can be constructed by starting from ten dimensions and assuming that six dimensions are compactified on a torus or another manifold. The geometrical parameters that characterise compact dimensions are often arbitrary. In the effective field theory, this is reflected by the presence of massless fields, the moduli, with the VEVs corresponding to six-dimensional radii, angles *etc.* that remain undetermined at the classical level due to flat directions of the scalar potential. In addition to the massless modes, a typical spectrum also contains the towers of Kaluza-Klein excitations with the masses quantised in units of inverse radii $\sim 1/R$.

The simplest and in some sense unique mechanism for “stringy” SUSY breaking at an arbitrary scale is by twisting the compact tori, *i.e.* by imposing a special type of boundary conditions in compact dimensions [11, 12]. A typical twist cuts out every second state of each Kaluza-Klein tower and eliminates the massless gravitino together with half of its tower. The remaining half of the gravitino tower starts with a massive spin 3/2 particle which can be identified as the gravitino of spontaneously broken supergravity with $m_{3/2} \sim 1/R$. In this way, the SUSY breaking scale $M_S \sim 1/R$ becomes tied up with a compact radius. From the supergravity point of view, twisted tori give rise to stringy realisations of no-scale models with vanishing potentials and zero cosmological constant at the tree level. SUSY breaking is due to a VEV of the auxiliary F -component of a supermultiplet containing the modulus T whose VEV determines R ; the modulino plays the role of the goldstino. As mentioned before, $\langle 0|T|0 \rangle = R$ remains arbitrary at the tree level. This flatness of the potential is due to a special moduli-dependence of the Kähler potential that follows directly from superstring theory. Furthermore, the loop corrections have no ultraviolet divergences since the string mass scale M provides a physical cutoff.

The usual pattern of soft-breaking terms induced by twisting is such that the scalar partners of quarks and leptons remain massless at the tree level whereas gauginos receive a common mass $m_{1/2} = m_{3/2}$ [12, 13]. Once supersymmetry is broken, the radiative corrections lift the flatness of the potential by generating a non-trivial potential for T . Minimisation of this potential with respect to T and to the Higgs field will fix their VEVs. The new VEV scale is defined as the energy at which the mass squared of the Higgs becomes negative and the breaking of electroweak symmetry occurs, and is expected to be $\sim M_P e^{-1/Y^2(M_P)}$, where Y is some Yukawa coupling. In this way, $M_S \sim 1/R$ can be hierarchically smaller than M_P provided that h is not too large.

A low supersymmetry breaking scale $M_S \sim 1$ TeV corresponds to a large internal dimension. A completely new, higher-dimensional world opens up above 1 TeV. From the four-dimensional point of view, the extra dimensions would manifest themselves by the presence of infinite towers of Kaluza-Klein excitations. Naïvely, this would seem to contradict superstring unification at 10^{17} GeV which is based on the logarithmic running of gauge coupling constants with the assumption of a desert between M_S and the unification mass. The reason why there is no contradiction is that Kaluza-Klein states are organised in multiplets of $N=4$ supersymmetry. An $N=4$ multiplet contains one vector boson, four two-component spinors and six real scalars. This leads to cancellation of the large radiative corrections among particles of different spin

and the evolution of gauge couplings remains logarithmic, as in four-dimensional theory, up to the Planck scale [14].

The perspective of probing extra dimensions at future colliders seems very appealing. Among the various Kaluza-Klein excitations of different spin, the easiest to detect are the vectors with the quantum numbers of the electroweak bosons. They would decay into quarks, leptons or into their SUSY partners; the lifetime can be estimated to be of order 10^{-26} seconds [13].

There exists another way of SUSY breaking which employs extra dimensions. A constant magnetic field, associated with a $U(1)$ gauge group, which points in the direction of extra dimensions, generates mass splittings within SUSY multiplets carrying non-zero $U(1)$ charges [15]. Here again, $M_S \sim 1/R$. The main difference between twisted and magnetised tori is that in the latter case a non-zero potential, and a possible electroweak symmetry breaking, are present already at the tree level.

To summarise, twisted and magnetised tori provide viable mechanisms for low-energy SUSY breaking in superstring theory. From the theoretical point of view the most important problem that requires further clarification is the string description of the vacuum rearrangement that leads to electroweak symmetry breaking and to the determination of M_S . For instance, at the string level, a non-vanishing one-loop cosmological constant leads to infinite tadpoles at two loops, therefore a consistent prescription for handling these divergences is necessary in order to obtain definite predictions for the soft-breaking terms.

4.2. Gaugino Condensation

Four-dimensional superstring theories usually contain very rich spectra that include not only the standard model sector but also hidden sectors which are very often associated with a whole new non-abelian gauge group. Dynamical supersymmetry breaking may then occur as a non-perturbative effect of hidden gauge interactions, much like in the supercolour idea mentioned before, and may be communicated to the observable sector via higher-dimensional interactions. Assuming that non-perturbative effects take place at energies much lower than $M \sim M_P$, they can be described within the framework of the effective field theory. This approach can only be justified *a posteriori*: once supersymmetry is found broken at $M_S \ll M$, one should argue that the respective physical mechanism remains unaffected by high-energy string physics.

As an example of a simplest hidden gauge sector, consider an asymptotically free QFT defined by a pure supersymmetric Yang-Mills system with an arbitrary gauge group. Non-perturbative dynamics of this theory have been studied extensively in the past in the context of global supersymmetry. In particular, there is a mass gap, and the lightest fermion, which is expected to be the superpartner of the glueball, has a mass of order of the strong interaction scale Λ [16]. Since there is no goldstino available, supersymmetry remains unbroken even at the non-perturbative level, as confirmed by Witten index theorem [2]. On the other hand, a non-perturbative effect that does certainly occur is the gaugino condensation [16] which gives rise to

$$\langle 0 | \lambda \lambda | 0 \rangle \sim \Lambda^3 \sim \mu^3 \exp\left(-\frac{3}{2\beta_0 g^2(\mu)}\right) \quad (13)$$

where μ is the renormalisation scale, $g(\mu)$ is the gauge coupling, and β_0 is the one-loop beta function coefficient.

There is an important difference between globally supersymmetric gauge theories and the effective field theories describing gauge interactions in superstring theory. In the latter case, the gauge couplings, similarly to other physical quantities, correspond to dynamical parameters which are determined by VEVs of some scalar fields. In heterotic superstring theory, a typical

gauge function which determines the gauge coupling at the string scale has the form

$$f_a(\phi) = S + f_a^{(1)}(T) = \frac{1}{g_a^2(S, T)} , \quad (14)$$

where the tree-level contribution depends universally on the dilaton S while the one-loop threshold corrections $f_a^{(1)}$ depend on the moduli T [17, 18]. As a result, the auxiliary field equations receive additional terms involving gaugino bilinears:

$$\bar{F}_{\phi} = \bar{F}_{\phi}(\text{BOSONS}) + \kappa^2 \left(\frac{\partial^2 \widehat{K}}{\partial \phi \partial \bar{\phi}} \right)^{-1} \frac{\partial(\frac{1}{g^2(\phi)})}{\partial \phi} \langle 0 | \lambda \lambda | 0 \rangle + \dots , \quad (15)$$

where $F_{\phi}(\text{BOSONS})$ is the bosonic part given by eq.(11). In this way, gaugino condensation breaks supersymmetry at $M_S \sim \Lambda^3/M_P^2$ [19]. The missing goldstino is found as a combination of the dilatino and the modulinos, as seen from eqs.(14) and (15).

The values of gauge couplings at the string scale, hence Λ and M_S , are all determined by the dilaton and moduli VEVs. In order to compute these VEVs one has to determine first the effective potential induced by non-perturbative effects. This can be done by integrating out the gauge degrees of freedom in the effective theory describing a coupled Yang-Mills – dilaton/moduli system [20]. The final result is the effective superpotential

$$\widehat{W}(S, T) \sim \Lambda^3 \sim M^3 \exp\left(-\frac{3}{2\beta_0 g^2(S, T)}\right) . \quad (16)$$

The moduli-dependence of \widehat{W} and of the respective potential is due to the one-loop threshold corrections $f_a^{(1)}(T)$, eq.(14). The form of these functions is well known, however there is no need to go into details to point out some basic features of the potential. The strongest constraint comes from the invariance of superstring theory under duality transformations $R \rightarrow 1/(RM^2)$ relating large and small radius compactifications. This duality is due to a complete symmetry between Kaluza-Klein excitations and string winding modes. It is reflected in the effective field theory, hence in the scalar potential, as a symmetry under modular transformations $T \rightarrow 1/T$. A potential symmetric under such a transformation has an obvious stationary point at the self-dual point $T = 1$. A more detailed analysis, using the explicit expressions for threshold corrections, shows that this corresponds to a minimum or that a true minimum with respect to T is located in the neighbourhood of the self-dual point. In this way, the radii are stabilised at a typical value $R \sim 1/M$.

On the other hand, the minimisation of the potential with respect to the dilaton S presents a more difficult problem. From the dilaton-dependence of gauge couplings, eq.(14), it follows that $\widehat{W} \sim \exp(-3S/2\beta_0)$. The respective potential falls off exponentially at large S and there is no stable minimum. There is of course a “runaway” vacuum at $S \rightarrow \infty$ corresponding to $\Lambda \rightarrow 0$ and unbroken supersymmetry. It is not surprising that the theory prefers to relax in a zero energy supersymmetric vacuum. It is very difficult to understand how a stable vacuum can exist at finite S . The formula (14) which is responsible for the exponential suppression of the superpotential is correct to all orders of perturbation theory [18]. A different dilaton dependence of gauge couplings, induced by some truly non-perturbative superstring effects, could in principle alter eq.(14) [21, 22]. However, a low M_S requires gaugino condensation to occur at $\Lambda \ll M$, and it is hard to imagine how genuinely superstring effects could interfere at such a low scale. The onset of these effects can be seen in the effective field theory as the appearance of interactions described by higher-derivative supergravity, but all of them are suppressed by the powers of Λ/M .

To summarise, there is a serious self-consistency problem with QFT description of SUSY breaking by gaugino condensation.^c There is much further work needed in order to provide superstring-theoretical description of non-perturbative QFT physics. From this point of view, recent developments in dualities and other non-perturbative aspects of superstring theory look very promising and go straight in the right direction.

5. New Results and Perspectives

Up to this point, there have not been many new results reported in this review. In the past year, most of field-theoretical studies of SUSY breaking have focussed on the following topics:

- model-building with dynamical SUSY breaking [5]
- general analysis of soft-breaking terms in the effective supergravity theory [23]
- studies of the effective actions describing gaugino condensation [24, 25]
- mass generation for the universal axion [25]
- strong-weak coupling duality-inspired dilaton stabilisation [22].

Recently, there have been many exciting new developments in superstring theory that bear excellent prognosis for a deeper understanding of SUSY breaking. Many mysterious “dualities” [26] have been discovered which allow exact determination of some physical quantities in $N = 2$ and $N = 1$ supersymmetric models. For instance, a $N = 2$ prepotential which usually contains perturbative and non-perturbative contributions can be computed in some cases exactly as a purely classical quantity in the dual theory [27]. All dualities known so far relate theories with equal number of supersymmetries, so they are not useful for SUSY breaking. There is no reason however why dual descriptions should not exist for $N = 1$ superstrings with Yang-Mills sectors that break supersymmetry by gaugino condensation. It would not be surprising if the dual descriptions involved twisted or magnetised tori; the two previous subsections might in fact describe different aspects of the same physical mechanism.

In summary, there is a clear advantage gained by promoting supersymmetry to a local symmetry: all known interactions can be described in one unified framework of supergravity. In supergravity models, SUSY breaking is transmitted from the hidden sector to the observable sector in a very natural way. Superstrings take us much farther, by offering a completely calculable framework with a physical short-distance cutoff. Many important aspects of SUSY breaking in superstring theory have already been understood. It remains however to put several pieces together to obtain a fully consistent picture; most likely, it will include some sort of superstring – nonsupersymmetric string dualities. In this way, superstring theory may finally offer some firm predictions that can be tested at future colliders.

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^cThis problem may be absent though in some models, with gauge groups consisting of several non-abelian factors *etc.*

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